Unparticle physics on direct CP violation

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The effects of the peculiar CP conserving phases in unparticle propagators are explored. We find that the phases have a great impact on CP-violation. We adopt the decays $B_d \to \pi^+ \pi^-$ and $B_d \to \ell^- \ell^+$ as the illustrators to demonstrate the influences of these phases on the direct CP asymmetries. In particular, we emphasize that unparticle physics is the only model suggested to date that could give the direct CP asymmetries in $B_d \to \ell^- \ell^+$ as large as 15%. We also point out that the unparticle phases could be probed in $B \to K^* \ell^+ \ell^-$ decays by using T-odd correlations.

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Recently, it has been argued in Refs. [1, 2] that a scale invariant sector with scale dimension $d_{\mathcal{U}}$ is associated with "unparticle stuff" which looks like a non-integral number $d_{\mathcal{U}}$ of invisible particles. The production of unparticles might be detectable by measuring the missing energy and momentum distributions in various processes [1, 2]. The unparticle physics phenomenology is further implemented in Refs. [3, 4]. Although it is still unclear how far the unparticle physics can be carried out, there exist many possible interesting experimental tests on unparticles. In particular, it has been illustrated in Ref. [2] that the peculiar CP conserving phases associated with the unparticle propagators in the time-like region lead to unusual CP conserving interference patterns between the time-like unparticle exchange amplitudes and the standard model (SM) amplitudes in $e^+ e^- \to \mu^+ \mu^-$. The effect of the virtual unparticle propagation has also been considered in Ref. [3]. It should be interesting to ask if there are other odd effects due to this type of the phases. In this study, we will examine the effects of these phases on CP violation.

In a decay process, the direct CP asymmetry (CPA) is defined by

$$\mathcal{A}_{CP} \equiv \frac{\bar{\Gamma} - \Gamma}{\bar{\Gamma} + \Gamma} \qquad (1)$$

where Γ ($\bar{\Gamma}$) is the partial decay rate of the (CP-conjugate) process. It is well known that

$$\mathcal{A}_{CP} \propto \sin\theta_w \sin\theta_{st}, \qquad (2)$$

where θ_w and θ_{st}, are CP violating (CPV) and CP conserving (CPC) phases, so called weak and strong phases, respectively. In the SM, the weak CPV phase is the unique phase in the 3×3 Cabibbo-Kobayashi-Maskawa (CKM) matrix [5], which appears in the CKM matrix elements of V_{td} and V_{ub} [6], expressed by $V_{td} = |V_{td}|e^{-i\beta}$ and $V_{ub} = |V_{ub}|e^{-i\gamma}$ [7]. With 10^8 $B\bar{B}$ pairs produced at B factories, the measurement on $\sin 2\beta$ through the Golden mode of $B_d \to J/\Psi K_s$ is determined to be 0.73 ± 0.04 [7]. It has no doubt that the mixing induced CPA in the B system is dominated by the CKM phase without the need of a strong phase.

To probe the phase effects, one can also use some T-odd correlations such as the well known triple spin-momentum product correlation in a three-body decay [8]. As the T-odd effects of the correlations are proportional to $\sin\theta_w \cos\theta_{st}$ as well as $\cos\theta_w \sin\theta_{st}$, they can be generated by either CPV or CPC phase. As a result, if the CPV phase in the process vanishes, i.e. $\theta_w = 0$, one can utilize the T-odd correlations to test the CPC phase of θ_{st}.

In this Letter, we will use B decay processes to illustrate the impact of the CPC phases in the unparticle propagators on CP violation. Explicitly, we will show that without the strong QCD phases, the large direct CPA in $B_d \to \pi^- \pi^+$ could be generated when the unparticle physics is introduced. We will also demonstrate that the unparticle phase could create non-zero direct CPAs in $B_d \to \ell^- \ell^+$ decays, which are vanishing small in most of other CP violating models due to the lack of strong phases. Furthermore, we consider T-odd observables in the decays of $B \to K^* \ell^+ \ell^-$ ($\ell = e, \mu$) to reveal the unparticle phases since there are T-odd effects due to the zero CPV weak phase for the $b \to s$ transition in the SM.

We start with the propagator of a scalar (vector) unparticle, given by [1, 2]

$$\int d^4x e^{ip \cdot x} \langle 0|T\left(O_{\mathcal{U}}^{(\mu)}(x)O_{\mathcal{U}}^{(\nu)}(0)\right)|0\rangle = i\Delta_{\mathcal{U}}^{S(V)}(p^2)\, e^{-i\phi_{\mathcal{U}}} \quad (3)$$

with

$$\Delta_{\mathcal{U}}^S(p^2) = \frac{A_{d_{\mathcal{U}}}}{2\sin(d_{\mathcal{U}}\pi)}\frac{1}{(p^2 + i\epsilon)^{2-d_{\mathcal{U}}}},$$

$$\Delta_{\mathcal{U}}^V(p^2) = \frac{A_{d_{\mathcal{U}}}}{2\sin(d_{\mathcal{U}}\pi)}\frac{-g^{\mu\nu} + p^\mu p^\nu/p^2}{(p^2 + i\epsilon)^{2-d_{\mathcal{U}}}}, \qquad (4)$$

where $\phi_{\mathcal{U}} = (d_{\mathcal{U}} - 2)\pi$, $O_{\mathcal{U}}^{(\mu)}$ is the scalar (vector) unpar-

ticle operator and

$$A_{d_{\mathcal{U}}} = \frac{16\pi^{5/2}}{(2\pi)^{2d_{\mathcal{U}}}} \frac{\Gamma(d_{\mathcal{U}} + 1/2)}{\Gamma(d_{\mathcal{U}} - 1)\Gamma(2d_{\mathcal{U}})} . \quad (5)$$

Note that in Eq. (3) the phase factor arises from $(-1)^{d_{\mathcal{U}}-2} = e^{-i\pi(d_{\mathcal{U}}-2)}$ and the vector operator is assumed to satisfy the transverse condition $\partial_\mu O_{\mathcal{U}}^\mu = 0$.

To study the unparticle physics, we write the effective interactions for unparticle operators coupling to leptons and quarks to be [1, 2]

$$\frac{c_A^\ell}{\Lambda_{\mathcal{U}}^{d_{\mathcal{U}}-1}} \bar{\ell}\gamma_\mu\gamma_5\ell O_{\mathcal{U}}^\mu + \frac{c_P^\ell}{\Lambda_{\mathcal{U}}^{d_{\mathcal{U}}}} \bar{\ell}\gamma_\mu\gamma_5\ell\partial^\mu O_{\mathcal{U}} +$$
$$\frac{c_{VA}^{q'q}}{\Lambda_{\mathcal{U}}^{d_{\mathcal{U}}-1}} (\bar{q}'q)_{V-A} O_{\mathcal{U}}^\mu + \frac{c_{SP}^{q'q}}{\Lambda_{\mathcal{U}}^{d_{\mathcal{U}}}} (\bar{q}'q)_{V-A}\partial_\mu O_{\mathcal{U}} \quad (6)$$

where $(\bar{q}'q)_{V-A} = \bar{q}'\gamma_\mu(1-\gamma_5)q$, c's are the dimensionless parameters and $\Lambda_{\mathcal{U}}$ is the energy scale below which the scale invariant unparticle fields emerge. For simplicity, in Eq. (6) we have only chosen the relevant interacting structures for $B_d \to \pi^-\pi^+$, $B_d \to \ell^-\ell^+$ and $B \to K^*\ell^+\ell^-$ and we have set that all c's are real numbers, as the general case is not the issue here.

In terms of the factorization approach, the decay amplitude for $B_d \to \pi^-\pi^+$ is expressed by

$$A(B_d \to \pi^-\pi^+) = \frac{G_F}{\sqrt{2}} V_{ub}^* V_{ud} a_1 m_B^2 f_\pi F_0^{B\pi}(m_\pi^2)$$
$$\times \left[1 + \chi_{\pi\pi} e^{-i\phi_{\mathcal{U}}} e^{-i\gamma}\right] \quad (7)$$

with

$$\chi_{\pi\pi} = \frac{8}{g^2 a_1 N_c} \frac{c_{VA}^{bd} c_{VA}^{uu}}{|V_{ub}^* V_{ud}|} \frac{A_{d_{\mathcal{U}}}}{2\sin d_{\mathcal{U}}\pi} \frac{m_W^2}{p^2} \left(\frac{p^2}{\Lambda_{\mathcal{U}}^2}\right)^{d_{\mathcal{U}}-1},$$

where $a_1 = C_2 + C_1/N_c$ is the effective Wilson coefficient [9] with N_c being the color number, $p^2 \sim m_B\bar{\Lambda}$ with $\bar{\Lambda} = m_B - m_b$ and f_π and $F_0^{B\pi}$ denote the pion decay constant and $B \to \pi$ form factor, respectively. We note that the scalar unparticle contributions have been ignored as they are suppressed by $m_B^2/\Lambda_{\mathcal{U}}^2$. Subsequently, the BR and direct CPA are found to be

$$\mathcal{B}(\pi^-\pi^+) = \mathcal{B}_0^{SM} \left(1 + \chi_{\pi\pi}^2 + 2\chi_{\pi\pi}\cos d_{\mathcal{U}}\pi\cos\gamma\right),$$
$$\mathcal{A}_{CP} = \frac{2\chi_{\pi\pi}\sin d_{\mathcal{U}}\pi\sin\gamma}{1 + \chi_{\pi\pi}^2 + 2\chi_{\pi\pi}\cos d_{\mathcal{U}}\pi\cos\gamma}, \quad (8)$$

where $\mathcal{B}_0^{SM} = \tau_{B_d} m_B^3 G_F^2 a_1^2 f_\pi^2 F_0^{B\pi}(m_\pi^2)|V_{ub}^* V_{ud}|^2/32\pi$ is the SM contribution.

The unknown parameter c_{VA}^{bd} can be directly constrained by the $B_d - \bar{B}_d$ mixing and the explicit relation is given by

$$\Delta M_{B_d}^{\mathcal{U}_V} \approx \frac{f_B^2}{m_B} \frac{A_{d_{\mathcal{U}}}}{2\sin d_{\mathcal{U}}\pi} \left(\frac{m_B^2}{\Lambda_{\mathcal{U}}^2}\right)^{d_{\mathcal{U}}-1} \left|c_{VA}^{bd}\right|^2 . \quad (9)$$

With $\Delta M_{B_d}^{\mathcal{U}_V} \leq 3 \times 10^{-13}$ GeV [7], $f_B = 0.2$ GeV, $d_{\mathcal{U}} = 3/2$ and $\Lambda_{\mathcal{U}} = 1$ TeV, the order of magnitude

for c_{VA}^{bd} could be known as $|c_{VA}^{bd}| \sim 10^{-4}$. The direct CPA in $B_d \to \pi^-\pi^+$ as a function of the scale dimension $d_{\mathcal{U}}$ with $\Lambda_{\mathcal{U}} = 1$ TeV is presented in Fig. 1, where the solid, dashed, and dash-dotted lines corresponds to $c_{VA}^{uu} = 0.05$, 0.1 and 0.5, respectively. The band in the figure denotes the world average 0.38 ± 0.07 [10]. From our results, we see clearly that the CPC phase in the unparticle physics provides a mechanism to produce a large CPA for $B_d \to \pi^-\pi^+$ decay. However, we remark that both experimental measurements [11] and theoretical results in and beyond the SM [12] are still inconclusive.
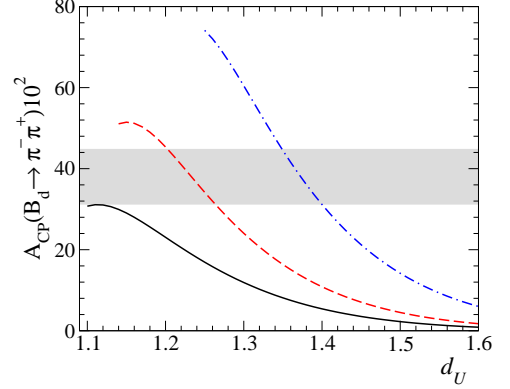

FIG. 1: Direct CP asymmetry in $B_d \to \pi^-\pi^+$ as a function of $d_{\mathcal{U}}$ with $\Lambda_{\mathcal{U}} = 1$ TeV, where the sold, dashed and dash-dotted lines correspond to $c_{VA}^{uu} = 0.05$, 0.1 and 0.5, respectively, and the band denotes the world average with 1σ errors.

Since the vector unparticle has only transverse degrees of freedom, it has no effects on $B_d \to \ell^-\ell^+$ ($\ell = e, \mu$ and τ). However, the scalar unparticle could give contributions to the decays. As a result, the decay amplitudes are found to be

$$A(B_d \to \ell^+\ell^-) = -i2m_\ell m_B^2 f_B \frac{c_{SP}^{bd}}{\Lambda_{\mathcal{U}}^{d_{\mathcal{U}}}} \frac{c_P^\ell}{\Lambda_{\mathcal{U}}^{d_{\mathcal{U}}}}$$
$$\times \Delta_{\mathcal{U}}^S(m_B^2) e^{-i\phi_{\mathcal{U}}} \bar{\ell}\gamma_5\ell; \quad (10)$$

and the decay BRs for $B_d \to \ell^-\ell^+$ are given by

$$\frac{1}{m_\ell^2}\mathcal{B}(B_d \to \ell^+\ell^-) = \kappa_B \left|Z_{SM}^B e^{-i\beta} + Z_{\mathcal{U}}^B e^{-i\phi_{\mathcal{U}}}\right|^2$$

where

$$\kappa_B = \frac{1}{m_\tau^2} \frac{\alpha^2 \mathcal{B}(B^+ \to \tau^+\nu_\tau)}{\pi^2 \sin^4\theta_W} \frac{\tau_{B_d}}{\tau_{B^+}},$$
$$Z_{\mathcal{U}}^B = \frac{16\pi\sin^2\theta_W}{g^2\alpha} \frac{c_{SP}^{bd} c_P^\ell}{|V_{ub}|} \frac{A_{d_{\mathcal{U}}}}{2\sin d_{\mathcal{U}}\pi} \frac{m_W^2}{m_B^2} \left(\frac{m_B^2}{\Lambda_{\mathcal{U}}^2}\right)^{d_{\mathcal{U}}},$$

and $Z_{SM}^B = |V_{tb}^* V_{td}| Y(x_t)/|V_{ub}|$ with $x_t = m_t^2/m_W^2$. Here, we have used the measured $B^+ \to \tau^+\nu_\tau$ decay to remove the uncertainty from f_B. Due to $m_t \gg m_W$, the function of $Y(x_t)$ is simplified as $Y(x_t) = 0.315\,x_t^{0.78}$ [9]. As the B meson is much heavier than leptons, we treat

the leptons as massless particles. The explicit m_ℓ comes from the equation of motion. In order to make the discussions independent of lepton species, we will concentrate on the quantity of $\mathcal{B}(B_d \to \ell^-\ell^+)/m_\ell^2$.

It is known that long-distance contributions which could introduce strong phases are negligible in the pure leptonic B decays. As a result, one cannot get sizable direct CPAs in $B_d \to \ell^-\ell^+$ decays in the SM although there exists a weak phase β. Moreover, no matter how many new weak CPV phases in new physics are injected, the conclusion cannot be changed until an intermediate state is involved, which carries an exotic CPC phase. Amazingly, the peculiar phase associated with the unparticle propagator is the unique phase that can generate non-zero CPAs in $B_d \to \ell^-\ell^+$.

Similar to $B_d \to \pi^-\pi^+$, c_{SP}^{bd} could be constrained by the $B_d - \bar{B}_d$ mixing, given by

$$\Delta M_{B_d}^{\mathcal{U}_S} \approx \frac{5}{3}\frac{f_B^2}{m_B}\frac{A_{d_\mathcal{U}}}{2\sin d_\mathcal{U}\pi}\left(\frac{m_B^2}{\Lambda_\mathcal{U}^2}\right)^{d_\mathcal{U}}\left|c_{SP}^{bd}\right|^2 , \quad (11)$$

where the matrix element $\langle B_d|\bar{b}(1-\gamma_5)d\,\bar{b}(1-\gamma_5)d|\bar{B}_d\rangle \approx -5/6f_B^2 m_B$ has been used. With $\mathcal{B}(B^+ \to \tau^+\nu_\tau) = 1.79 \times 10^{-4}$ [14] and $\Lambda_\mathcal{U} = 1$ TeV, we present our numerical results of the BR and the direct CPA in Fig. 2, where the solid, dashed and dash-dotted lines denote $c_P^\ell = 0.1, 0.5$ and 1.0, respectively. From Fig. 2a, we
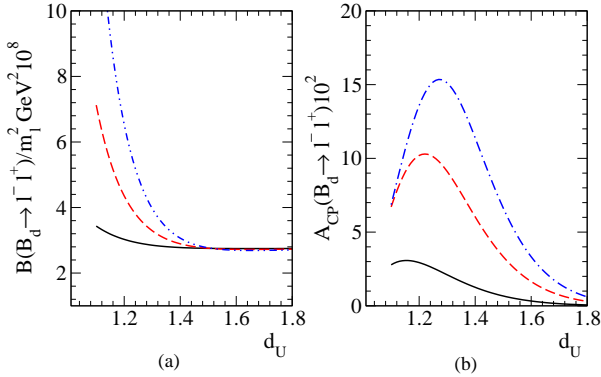

FIG. 2: (a) BR/m_ℓ^2 and (b) direct CP asymmetry in $B_d \to \ell^-\ell^+$ as functions of $d_\mathcal{U}$ with $\Lambda_\mathcal{U} = 1$ TeV, where the sold, dashed and dash-dotted lines correspond to $c_P^\ell = 0.1, 0.5$ and 1.0, respectively.

see that when the scale dimension $d_\mathcal{U}$ is less than 1.4, $\mathcal{B}(B_d \to \ell^-\ell^+)/m_\ell^2$ becomes very sensitive to it. The merging flat lines for $d_U > 1.4$ correspond to the SM contributions. As shown in Fig. 2b, it is very interesting to see that $\mathcal{A}_{CP}(B_d \to \ell^-\ell^+)$ can be as large as 15%. This is a very unique result in the unparticle physics, which is different from most of other new physics. If a large deviation in BR for $B_d \to \ell^-\ell^+$ is found experimentally, of course, it could be a signal for the existence of new physics; but, it could not tell us what kind of new physics involved. However, if a nonzero CPA in $B_d \to \ell^-\ell^+$ is observed, the unparticle physics definitely plays the most important role.

Finally, we study the interesting T-odd operators in $B \to K^*\ell^+\ell^+$ with polarized K^* [13]. With the axial-vector interactions in Eq. (6), the corresponding effective Hamiltonian for $b \to s\ell^+\ell^-$ with the SM effects is found to be

$$\begin{aligned}\mathcal{H}_{\text{eff}} &= \frac{G_F\alpha_{em}V_{ts}V_{tb}^*}{\sqrt{2}\pi}\Big[\big(C_9^{\text{eff}}(\mu)\bar{s}\gamma_\mu P_L b \\ &\quad -\frac{2m_b}{q^2}C_7(\mu)\bar{s}i\sigma_{\mu\nu}q^\nu P_R b\big)\bar{\ell}\gamma^\mu\ell \\ &\quad +C_{10}^\mathcal{U}\bar{s}\gamma_\mu P_L b\,\bar{\ell}\gamma^\mu\gamma_5\ell\Big] ,\end{aligned} \quad (12)$$

where C_9^{eff} and C_7 are the Wilson coefficients of the SM [9]. We note that C_9^{eff} includes the $c\bar{c}$ resonant effects. The unparticle contribution only appears in $C_{10}^\mathcal{U}$ and is given by

$$C_{10}^\mathcal{U} = C_{10} + \frac{16\pi c_{VA}^{bs}c_A^\ell}{g^2\alpha_{em}V_{ts}V_{tb}^*}\frac{m_W^2}{q^2}\left(\frac{q^2}{\Lambda_\mathcal{U}^2}\right)^{d_\mathcal{U}-1} . \quad (13)$$

To study the T-odd spin-momentum correlation effects, we need to consider the K^* polarizations, i.e., we have to consider the decay chain $B \to K^*\ell^+\ell^- \to K\pi\ell^+\ell^-$. Hence, in terms of the obtained Hamiltonian, the differential decay rate for $B \to K\pi\ell^+\ell^-$ is obtain by

$$\begin{aligned}&\frac{d\Gamma}{d\cos\theta_K d\cos\theta_\ell d\phi dq^2} = \frac{3\alpha_{em}^2 G_F^2 |V_{ts}V_{tb}^*|^2 |\vec{p}|}{2^{14}\pi^6 m_B^2}\\ &\times Br(K^* \to K\pi)\{\cdots + 2\sin 2\theta_K \sin\theta_\ell \sin\phi \\ &\times\big(Im\mathcal{M}_1^0(\mathcal{M}_2^{+*} + \mathcal{M}_2^{-*}) - Im(\mathcal{M}_1^+ + \mathcal{M}_1^-)\mathcal{M}_2^{0*})\} ,\end{aligned} \quad (14)$$

where for simplicity we just display the terms associated with dominant T-odd effect and the irrelevant terms are denoted by $\cdots$, $\theta_{\ell(K)}$ is the polar angles of $\ell^-(K)$ in the rest frame of $\ell^+\ell^-(K^*)$, ϕ represents the relative angle of the decaying plane between $K\pi$ and $\ell^+\ell^-$, q^2 is the invariant mass of $\ell^+\ell^-$ and $|\vec{p}| = [((m_B^2 + m_{K^*}^2 - q^2)/(2m_b))^2 - m_{K^*}^2]^{1/2}$. In Eq. (14), $\mathcal{M}_i^0$ and $\mathcal{M}_i^\pm$ denote the longitudinal and transverse polarizations of K^* and their explicit expressions are given by

$$\mathcal{M}_{a\mu}^{(\lambda)} = if_1\varepsilon_{\mu\nu\alpha\beta}\epsilon^{*\nu}(\lambda)P^\alpha q^\beta + f_2\epsilon_\mu^*(\lambda) + f_3\epsilon^*(\lambda)\cdot qP_\mu,$$

respectively, where $P = p_B + p_{K^*}$, $q = p_B - p_{K^*}$ and $a = 1[2]$ while $f_i = h_i [g_i]$ $(i = 1, 2, 3)$ with

$$\begin{aligned}h_1 &= C_9^{\text{eff}}(\mu)V(q^2) - \frac{2m_b}{q^2}C_7(\mu)T(q^2),\\ h_{2(3)} &= -C_9^{\text{eff}}(\mu)A_{0(1)}(q^2) + \frac{2m_b}{q^2}C_7(\mu)T_0(q^2),\\ g_1 &= C_{10}^\mathcal{U}V(q^2), \quad g_{2(3)} = -C_{10}^\mathcal{U}A_{0(1)}(q^2),\end{aligned}$$

and $V(q^2)$, $A_{0(1)}(q^2)$ and $T_{0(1)}$ being the $B \to K^*$ form factors [13].

To illustrate the T-odd effects due to unparticles, we concentrate on the T-odd observable defined by $\langle\mathcal{O}_T\rangle =$

$\int \mathcal{O}_T d\Gamma$ [13] where $\mathcal{O}_T$ is a T-odd five-momentum correlation, given by

$$\mathcal{O}_T = \frac{(\vec{p}_B \cdot \vec{p}_K)\,(\vec{p}_B \cdot (\vec{p}_K \times \vec{p}_{\ell^+}))}{|\vec{p}_B|^2\,|\vec{p}_K|^2\,\omega_{\ell^+}}$$

with $\omega_{\ell^+} = q \cdot p_{\ell^+}/\sqrt{q^2}$. In the K^* rest frame, we note that $\mathcal{O}_T = \cos\theta_K \sin\theta_K \sin\theta_\ell \sin\phi$. The statistical significance of the observable can be determined by

$$A_T(q^2) = \frac{\int \mathcal{O}_T d\Gamma}{\sqrt{(\int d\Gamma)(\int \mathcal{O}_T^2 d\Gamma)}}\,. \qquad (15)$$

With the constraint of Δm_{B_s} [7] on c_{VA}^{bs} and $c_A^\ell = 5 \times 10^{-4}$, the results with unparticles on the T-odd observable are displayed in Fig. 3. We note that the SM
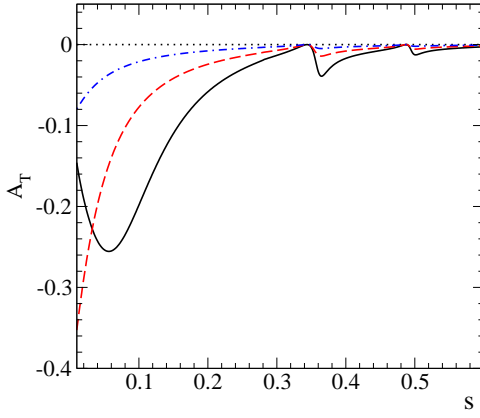

FIG. 3: Statistical significance A_T of the T-odd observable in $B_d \to K^{*0}\ell^+\ell^-$ as a function of $s = q^2/m_B^2$, where the solid, dashed and dot-dashed lines correspond to $d_{\mathcal{U}} = 1.2$, 1.4 and 1.6 while the dotted line is the SM result, respectively.

contribution vanishes as the associated CKM matrix elements of V_{ts} and V_{tb} contain no CPV phase. Moreover, it is interesting to point out that even if we include the resonance effects, the T-odd effects also get canceled in the SM. As seen from Fig. 3, the unparticle effects could make the statistical significance as large as 20%. It is clear if there is no any direct CP violating effect in the $b \to s$ transition, e.g., no $\sin 2\beta$ difference between $B_d \to J/\Psi K_S$ and $B_d \to \phi K_S$, any signal found in the T-odd obervables of $B \to K^*\ell^+\ell^-$ will directly reveal unparticle phase effects.

In summary, we have illustrated that the peculiar CPC phases in the unparticle propagators can play very important roles on the direct CPAs in $B_d \to \pi^-\pi^+$ and $B_d \to \ell^-\ell^+$ and the T-odd observables in $B \to K^*\ell^+\ell^-$. We have demonstrated that unparticle physics is the only model suggested to date that could give $\mathcal{A}_{CP}(B_d \to \ell^-\ell^+)$ as large as 15%. In addition, we have considered the T-odd effects due to the CPC unparticle phase in $B \to K^*\ell^+\ell^-$, which vanish in the SM as the CPV weak phase is negligible in the $b \to s$ transition. Finally, we remark that the direct CPA for $B_d \to \tau^+\tau^-$ could be accessible at a future super-B factory, such as the SuperKEKB which could produce $O(10^{10})$ $B\bar{B}$ pairs in its initial stage [15]. Moreover, 4400 events/year for $B \to K^*\ell^+\ell^-$ decays will be produced at the LHCb, corresponding to the accuracy of the rate asymmetries being around percent level [16]. Clearly, we have a great chance to observe the virtual unparticle phase effects in the B system.

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